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Contributions to the Cosmic Ray Flux above the Ankle: Clusters of Galaxies

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ABSTRACT

Motivated by the suggestion of Kang, Ryu & Jones (1996) that particles can be accelerated to high energies via diffusive shock acceleration process at the accretion shocks formed by the infalling flow toward the clusters of galaxies, we have calculated the expected particle flux from a cosmological ensemble of clusters. We use the observed temperature distribution of local clusters and assume a simple power-law evolutionary model for the comoving density of the clusters. The shock parameters such as the shock radius and velocity are deduced from the ICM temperature using the self-similar solutions for secondary infall onto the clusters. The magnetic field strength is assume to be in equipartition with the postshock thermal energy behind the accretion shock. We also assume that the injected energy spectrum is a power-law with the exponential cutoff at the maximum energy which is calculated from the condition that the energy gain rate for diffusive shock acceleration is balanced by the loss rate due to the interactions with the cosmic background radiation. In contrast to the earlier paper we have adopted here the description of the cosmic ray diffusion by Jokipii (1987) which leads to considerably higher particle energies. Finally the injected particle spectrum at the clusters is integrated over the cosmological distance to earth by considering the energy loss due to the interactions with the cosmic background radiation. Our calculations show that the expected spectrum of high-energy protons from the cosmological ensemble of the cluster accretion shocks could match well the observed cosmic ray spectrum near 10^{19}eV with reasonable parameters and models if about 10^{-4} of the infalling kinetic energy can be injected into the intergalactic space as the high energy particles.

Key words: Cosmic Rays -Hydrodynamics -Particle Acceleration-clusters of galaxies

1 INTRODUCTION

It is widely believed that diffusive shock acceleration is the mechanism from which cosmic rays get their energy. There are various models to account for the origin of cosmic rays below about $3{\times}10^{18}\,\mathrm{eV} = 3\,\mathrm{EeV}$. At low energies, up to about 10^{14} eV, supernova explosions into the interstellar medium give a reasonable and successful explanation for the data for protons. At higher energies, several models have been proposed, most notably a galactic wind termination shock (Jokipii & Morfill 1987), and multiple shocks in an ensemble of OB superbubbles and young supernova remnants (Axford 1992). A comprehensive, albeit tentative, theory has been proposed by Biermann (1993 and later papers) that explains the cosmic ray spectrum, with its chemical abundances and the knee feature as resulting from a combination of supernova explosions into the interstellar medium, and supernova explosions into strong stellar winds of progenitor Wolf-Rayet stars (Biermann 1993; Biermann & Cassinelli 1993; Biermann 1995). Above the "ankle" at about 3 EeV (the ultra-high energy cosmic ray regime, called UHECR hereafter), a simultaneous change in spectrum and composition of the cosmic ray spectrum (Bird et al. 1994) suggests a change of origin. As argued already by Cocconi (1956) particles of higher energy need to come from outside our galaxy due to the very large Larmor radius, and therefore need to be accelerated in extragalactic sources. Here we concentrate on these high energy particles that almost certainly come from extragalactic sources.

It is well known that the extragalactic cosmic ray spectrum must show the Greisen-Zatsepin-Kuzmin (GZK) cutoff at about 60 EeV due to interactions with the cosmic background radiation (CBR), regardless whether protons, heavy nuclei or photons are considered as the energetic particles (Greisen 1966; Zatsepin & Kuzmin 1966). Pair production and the cosmological evolution of

radiation backgrounds set limits for the distance of cosmic rays even at lower energies, thus any extragalactic model of cosmic ray origin must propose at least some sources in our cosmological neighborhood which can account for the highest energies. In an expansion of an earlier proposition by Biermann & Strittmatter (1987) which predicted maximum proton energies in active galaxies of 10^{21} eV, Rachen & Biermann (1993) developed a model to accelerate cosmic rays up to a few 100 EeV at strong shocks at the end of extended jets in powerful radio galaxies (FR-II radio galaxies, Fanaroff & Riley 1974). It has been shown that even for a moderate proton content in the jets this model can generally account for the cosmic ray flux above the ankle, and is consistent with air shower data suggesting a takeover from heavy to light nuclei in this energy range (Rachen, Stanev & Biermann 1993). It can be shown that this model can be extended to a maximum particle energy in the source of about 4×10^{21} eV (Biermann 1996). One prediction of this particular model has been tested successfully, and that is the expected correlation of arrival directions of very high energy cosmic rays with the large scale distribution of radio galaxies, in the supergalactic plane (Stanev et al. 1995; Hayashida et al. 1996).

However, the uncertain proton content of the jets, the energy limitation for cosmic rays due to the relatively small acceleration region in hot spots, the finite life time of hot spots, and the large distance of the closest well known FR-II radio galaxy set strong constraints on the predictability of the model, and leaves room for other contributions to the highest energy cosmic rays. In particular, after the detection of a 320 EeV air shower by Fly's Eye (Bird et al. 1994), and another 200 EeV event by the AGASA ground array (Hayashida et al. 1994), the various acceleration models have been critically reviewed (e.g. Elbert & Sommers 1995; Biermann 1996) and new models have been suggested. Most of them are based on astrophysical objects whose physical properties are under controversial discussion, as the decay of topological defects (Bhattacharjee 1991; Sigl, Schramm & Bhattacharjee 1994; Protheroe & Johnson 1996; Protheroe & Stanev 1996), or rapid acceleration in Gamma Ray Burst sources (Milgrom & Usov 1995; Vietri 1995; Waxman 1995), but the poor statistics of events above 100 EeV allows a variety of other explanations, including the earlier proposal of radio galaxy origin (Rachen 1995; Stanev et al. 1995; Biermann 1995, 1996).

Recent work shows that there could be sites for shock acceleration in extragalactic space alternative to the strongly constrained acceleration in radio galaxy hot spots. According to hydrodynamic simulations of large scale structure formation (e.g. Kang et al. 1994a; Cen & Ostriker 1994), accretion shocks are formed in the baryonic component around non-linear structures collapsed from the primordial density inhomogeneities as a result of gravitational instability. Those structures can be identified as pancake-like supergalactic planes, still denser filaments, and clusters of galaxies which form at intersections of pancakes, in any variants of the many cosmological models. They are surrounded by the hot gas heated by the accretion shocks and the particles can be accelerated to very high energies at these shocks via first order Fermi process. Kang, Jones & Ryu (1995) and Kang, Ryu & Jones (1996, KRJ96 hereafter) suggested that the accretion shocks around the clusters of galaxies could be as fast as 1000–3000 $\mathrm{km\,s^{-1}}$ and so could be good acceleration sites for the UHECRs up to several 10 EeV, provided there is a turbulent magnetic field so that the diffusion is in the Bohm limit, and if the magnetic field around the clusters is order of microgauss. We note here the maximum energy could be shifted to super-GZK values, if the field geometry is close to being quasi-perpendicular (Jokipii 1987). The significance of the accretion shocks is that they are the largest and longest lived shocks in the universe, so that they naturally get past the primary factors normally limiting Fermi acceleration to such high energies. Especially the accretion shocks formed around the clusters of galaxies which have the deepest gravitational potential well are the strongest and thus could accelerate the particles to the maximum possible energy. Independently Norman, Melrose & Achterberg (1995) showed with the help of a similar argument that the shocks associated with the large-scale structure formation could accelerate the protons up to $E_{\rm max} = 50$ EeV if there is a primordial field of 1–10 nanogauss, or if microgauss field can be self-generated in shocks.

In the present study we have estimated the contribution of the CR protons from an ensemble of the cluster accretion shocks distributed in the universe by adopting some simple models for the accretion flows onto clusters, the strength, geometry, degree of irregularities of the magnetic field near the cluster accretion shocks, and the cosmological evolution of the cluster distribution. The details of the models will be given in §2. In §3 the estimated CR proton spectrum has been compared with actual observations in order to see if the proposed origin could explain the energy spectrum of UHECRs with reasonable physical parameters and models. We discuss the implication of these results and models also in §3.

We adopt for the following $\Omega_0 = 1$ and write the Hubble constant as $H_0 = 100\,h\,\mathrm{km/s/Mpc}$; in numerical calculations we use $h = 0.75$.

2 MODELS

2.1 Accretion Shocks around Large Scale Structures

It is generally accepted that both galaxy distribution in the observed Universe and matter distribution in numerically simulated universe in most of generic cosmological scenarios show sheet-like and filamentary structures on large scale (e.g. de Lapparent, Geller & Huchra 1991; Melott 1987; White et al. 1987). Numerical simulations based on even a hierarchical clustering model such as variants of CDM models also indicate that the dominant non-linear structure is a network of filaments (White 1996) and the filaments became longer, straighter and clumpier as the hierarchical collapse proceeds (Summers 1996). Clusters of galaxies, especially the richest ones, form mostly at the vertices where several filaments intersect. They grow nonlinearly by gravitationally attracting matter along the filaments and also strengthen and align the filaments at the same time. The dominance of filaments in the visual impression of large scale matter distribution could be understood *in part* by the facts that the filaments as mostly intersecting curves between pancakes should naturally have higher density contrast than pancakes themselves (see West, Villumsen & Dekel 1991), and that a simulated universe is most often presented as a two-dimensional projection or a slice cut where sheetlike structures cannot be easily seen. It also depends to some degree on choosing a right level of density contrast. Thus, for example, knots (clusters) will become dominant if one chooses to look at only the highest density peaks.

Although formation of shocks and subsequent heating of the gas by them, when matter accretes toward these large-scale coherent structures, are implicit in any cosmological hydrodynamic simulations, the visual impression of accretion shocks has became clear only after high resolution Eulerian codes were introduced to the numerical cosmology (e.g. Ryu et al. 1993; Bryan et al. 1995). Interested readers are referred to Kang et al. (1994b) for the details of a comparison of various cosmological hydrodynamic

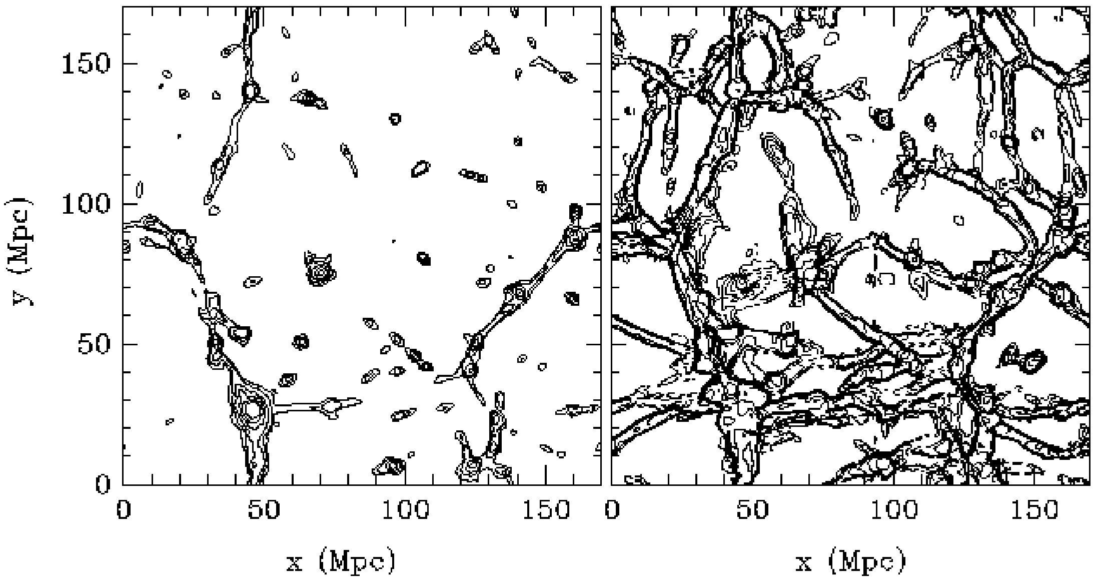

Figure 1. A representative slice cut of a simulated universe in a standard cold dark matter model. The left panel shows the X-ray luminosity distribution while the right panels shows the gas temperature distribution. See text for the details.

codes. They have demonstrated the accretion shocks can be captured clearly in the simulations done by these high-resolution Eulerian codes, while their existence is not so obvious in the simulations done by the particle codes based on the SPH method (see their Figs. 5). Fig. 1 shows a slice cut through a simulated universe based on the standard CDM model previously reported by Kang et al. (1994a). Readers are referred to their paper for the details of the cosmological model parameters, which should not be crucial to the current discussion. The left panel shows the X-ray luminosity modeled by $L_{\rm x} = \rho_{\rm b}^2 T^{1/2}$ and could be a representative distribution of X-ray clusters. Being weighted toward the higher density region, this shows the knot-like distribution with some alignments of several knots into filaments/sheets. The right panel shows the gas temperature distribution. The hotter region of $T \geq 10^6$ K is represented by solid contour curves, while the colder region of $T < 10^6$ K by dotted contour curves. Shocks can be seen clearly as strong gradients in temperature distribution, and they encompass the region of moderately overdense regions of sheetlike/filamentary structures. The most prominent cluster in the lower-left corner has been shown in Fig. 1 of KRJ96 in which the velocity pattern clearly shows accretion flows from the background IGM and along the filaments toward the cluster. These figures demonstrate that the accretion shocks associated with the large-scale coherent structures do indeed exist, even though any direct observations of such shocks have not been made yet. It is apparent from the diffusive shock acceleration theory that some particles are accelerated by these large-scale shocks. The accretion shocks around clusters of galaxies, being the fastest shocks generated by the deepest potential wells, are our focus in the present paper.

2.2 Self-Similar Evolution of Clusters

The temperature of the intracluster medium (ICM) within a cluster can be obtained with a reasonable accuracy from the X-ray observations (e.g. David et al. 1993). Fortunately, important physical parameters such as the velocity of an accretion shock, the velocity dispersion of galaxies, and the depth of the potential well of the cluster can be related rather well with the ICM temperature. Thus one can deduce such parameters from the observed temperature of X-ray clusters, for example, by assuming that the gas is in a hydrostatic equilibrium with the gravitational potential due to total mass including both the dark and baryonic matter. Here we take another approach in which one-dimensional (1D) spherical collapse of clusters is modeled as the collapse of an initially overdense point-mass perturbation followed by the secondary infall of background medium. Such models have been studied both semi-analytically (Fillmore & Goldreich 1984; Bertschinger 1985) and numerically (Ryu & Kang 1996). In such an accretion flow the infalling baryonic matter is stopped by a shock, while the collisionless matter forms many caustics. In this paper we assume that the evolution of clusters and the properties of the accretion shocks can be represented by such flows. Although this 1D approach does not account for the virialization of the central region of clusters, we will show below essential characteristics of real clusters can be related with parameters of the accretion shock through this model.

One can expect the evolution of such a flow and its accretion rate should be dependent upon the expansion of background universe for the given initial power spectrum of density perturbations (i.e. $P(k)$ is constant for a point-mass perturbation). The flow solution in an $\Omega_0 = 1$ universe approaches a self-similar form (Bertschinger 1985) due to the scale-free nature of the problem and so can be treated analytically. In a low density universe (i.e. $\Omega_0 < 1$), however, either open or flat with the cosmological

constant, the solution is not self-similar and can be studied only numerically (Ryu & Kang 1996). Here we will consider the accretion flows only in the Einstein-de Sitter universe ($\Omega_0 = 1$), and adopt the one-dimensional (1D), self-similar accretion given by Bertschinger (1985) as an evolutionary model for clusters and the accretion shocks.

Since the solution is self-similar, the shock parameters such as the radius and velocity of the accretion shock, and the post-shock gas temperature can be uniquely determined by a single parameter at a given epoch. For example, the mass contained inside the outermost caustic, $M_{\rm c} = M(r < r_{\rm c})$ at present epoch can be such a parameter. The radius and velocity of the accretion shock, and the postshock temperature at $r = 0.3 r_{\rm s} = 0.64 h^{-1}$ Mpc are given by $r_{\rm s} = 2.12 h^{-1} {\rm Mpc} (M_{\rm c}\ h/10^{15} M_\odot)^{1/3}$, $V_{\rm s} = 1.75{\times}10^3\ {\rm km\,s^{-1}} (M_{\rm c}\ h/10^{15} M_\odot)^{1/3}$, and $T_{\rm c}(0.3r_{\rm s}) = 6.06 {\rm keV} (M_{\rm c}\ h/10^{15} M_\odot)^{2/3}$, respectively (Ryu & Kang 1996). The reason that $T_{\rm c}(0.3r_{\rm s})$ is an interesting quantity will be given shortly. If one follows the evolution of a given perturbation, on the other hand, the length scale of the accretion flow which is proportional to the turn-around radius will grow with time as $r_{\rm ta} \propto t^{8/9} \propto (1+z)^{-4/3}$. The position of the shock, $r_{\rm s}$, in units of $r_{\rm ta}$ and the velocity of the shock, $V_{\rm s}$, in units of $r_{\rm ta}/t$ are fixed, so the shock radius and velocity for the same perturbation evolves as $r_{\rm s} \propto (1+z)^{-4/3}$ and $V_{\rm s} \propto (1+z)^{1/6}$.

Here we attempt to relate the physical parameters of the accretion shocks with observed temperature of X-ray clusters which often represents the emission weighted temperature in the core within $\sim 0.5 h^{-1}$ Mpc. We note that the temperature of the postshock gas in the 1D self-similar solution increases toward the center, in fact, to infinity, while observed ICM temperature distribution of most clusters is nearly isothermal, and in some cases shows a clear depression towards the center due to cooling. This discrepancy comes about, because the extrapolation of the self-similar evolution to $t \to t_i$ ($r \to 0$) is not valid during the initial collapse when the accreted mass is not much greater than the mass inside the initial perturbation. Also the 1D self-similar solution does not account for the virialization of the core region. According to the 3D cosmological hydrodynamic simulations (e.g. Crone, Evrard & Richstone 1994; Kang et al. 1994a; Navarro, Frenk & White 1995), the ICM is shock heated to the virial temperature and then settles into hydrostatic equilibrium with an approximate isothermal structure, which is consistent with the observed temperature distribution of the clusters. Cosmological SPH simulations (Navarro et al. 1995; Evrard, Metzler & Navarro 1995) showed that simulated clusters of different masses in fact have similar structures when scaled to a fixed density contrast (e.g. $\delta(r) = \bar\rho(r)/\rho_{\rm crit} \gtrsim 200{-}500$). Navarro et al. also showed that the temperature profile of simulated clusters can be approximated within a factor of two by that of the 1D self-similar solutions for the outer region, $r \gtrsim 0.3 r_{200}$ (where r_{200} is the radius at $\delta = 200$), while the inner region, $r < 0.3 r_{200}$, is isothermal. From this consideration, we have made a simple approximation that the observed X-ray temperature of a cluster is similar to $T_{\rm c}(0.3r_{\rm s})$ of the self-similar solutions, and so it provides the direct information about the shock velocity and radius from the self-similar solutions.

From the equations given above which relates $r_{\rm s}$, $T_{\rm c}(0.3r_{\rm s})$, and $V_{\rm s}$ with the cluster mass $M_{\rm c}$ and their redshift dependences, one can find them as a function of $kT_{\rm obs} = T_{\rm c}(0.3r_{\rm s})$ and z in an $\Omega_0 = 1$ universe as follows.

$$r_{\rm s} = 2.12 h^{-1}\,{\rm Mpc}\ \left(\frac{kT_{\rm obs}}{6.06{\rm keV}}\right)^{1/2} (1+z)^{-3/2} \qquad (1)$$

$$V_{\rm s} = 1.75{\times}10^3\ {\rm km\,s^{-1}} \left(\frac{kT_{\rm obs}}{6.06{\rm keV}}\right)^{1/2}. \qquad (2)$$

Evrard et al. , on the other hand, showed that r_{500} where $\delta(r_{500}) = 500$ is a *conservative* estimate for the boundary separating the inner virialized region from the outer infalling flow. They gave the scaling relation $r_{500} = 0.965\, h^{-1}\ {\rm Mpc} \times (kT/6.06\ {\rm keV})^{1/2}$ for clusters at $z = 0$ for an $\Omega_0 = 1$ universe. The fact that the ratio of the shock radius of the self-similar solution to the characteristic radius of virialized region in 3D simulation is $r_{\rm s}/r_{500} = 2.2$ seems reasonable, since $r_{\rm s}$ corresponds to the radial position at $\delta \sim 80$ and so it should be larger than r_{500}. According to Ryu & Kang (1996), the clusters of a given temperature have smaller accretion velocity by less 20% for $\Omega_0 = 0.3$ in either open or flat with $\Lambda_0 \neq 0$ universes. Since we use observed temperature distribution for the cluster abundance to be discussed below, the main results of our model should remain valid for a low Ω_0 cosmology.

According to Bertschinger (1985), the gas density upstream to the shock is $\rho_1 = 4.02 \Omega_{\rm b}\, \rho_{\rm crit}(z) = 7.56{\times}10^{-29} h^2\,{\rm g\,cm^{-3}} \times \Omega_{\rm b}(1+z)^3$. Thus all the shock parameters necessary for our model can be obtained from the observed redshift and X-ray temperature. In the following we adopt $\Omega_{\rm b} = 0.06$ in all numerical calculations.

2.3 Magnetic Field and Diffusion Coefficient

The strength and morphology of the intergalactic magnetic fields remain largely unknown and observational tasks to detect them are very challenging even with today's technology (see, for a review, Kronberg 1994). Theoretical study on the generation of the primordial fields and their subsequent amplification during the structure formation is also still in its infancy. In fact it is not clear at all that there was a primordial magnetic field. Recent study by Kulsrud et al. (1996), however, showed that a weak seed field can be generated at shocks and then amplified via the protogalactic turbulence during the structure formation up to equipartition with the turbulence. An alternative theory (Biermann 1996) uses magnetic stellar winds as the sources of magnetic fields for galaxies and their environment. On observational fronts, there exist some concrete observations that can be used in inferring the general distribution of the magnetic fields in intergalactic space. The current observational upper limit for a large-scale, pervading field is about $10^{-9}\, r_0^{-1/2}$ gauss (Kronberg 1994), where r_0 is the assumed reversal scale of the magnetic field topology in units of 1 Mpc. Thus, if the reversal scale throughout the universe were larger than 1 Mpc, such as $\approx 30 h^{-1}$ Mpc, the bubble-scale of the galaxy distribution, then this upper limit would be reduced by a factor of 6.3 (for $h = 0.75$). The magnetic fields inside typical galaxies are observed to be order of 3-10 μgauss, and magnetic fields near the μgauss level are also common in core regions of rich clusters according to many recent observations (Kim, Tribble & Kronberg 1991; Taylor & Perley 1993; Taylor, Barton & Ge 1994); some cooling flow clusters clearly show magnetic fields even higher than those typical in the interstellar medium of galaxies (see Kronberg 1994). Studies on the field generation via dynamos in the cooling flows (Ruzmaikin et al. 1989) and the field inputs from radio galaxies inside clusters (Enßlin et al. 1996) have attracted some attention recently. Here we are concerned most with the fields near the accretion shocks on larger scale than cluster core, that is, within $\sim 5 h^{-1}$ Mpc around clusters of galaxies. Observation of such fields has been attempted by Kim et al. (1989) in which fields of 0.1 μgauss level were deduced *in the plane of the supercluster* connecting the Coma cluster and A1367 by assuming equipartition between the magnetic fields and relativistic particles.

This "bridge" of emission region of $\sim 1.125h^{-1}$Mpc seems to be infalling toward the Coma and is within the supergalactic plane (so presumably inside the accretion shocks). Vallée (1990,1993), on the other hand, suggested the existence of a magnetic field component of $\sim 1.5\,\mu$gauss in a 10 Mpc region around the Virgo cluster, in this case using an assumed length scale of 10 Mpc. This might represent the fields outside the accretion shocks encompassing the supergalactic plane.

Here we adopt a simple, but common assumption that the ICM magnetic field is in rough equipartition with the thermal energy of the gas. Then the field strength in postshock region can be estimated from the thermal energy inside the shock (i.e. in the postshock region) according to

$$E_{\rm th} = 1.5 \frac{\rho_2}{\mu m_{\rm H}} k T_2 \sim \frac{B_2^2}{8\pi} \tag{3}$$

where $\rho_2(z) = 4\rho_1(z)$, $\mu = 0.61$, and $T_2 = 2.69{\times}10^7\,{\rm K} \times (kT_{\rm obs}/6.06{\rm keV})$. Then the postshock field is given by

$$B_2 = (1.71\,\mu{\rm gauss})\ f_B h \left[\frac{\Omega_{\rm b}}{0.06}\,\frac{kT}{6.06{\rm keV}}\right]^{\frac{1}{2}} \left[1+z\right]^{\frac{3}{2}}, \tag{4}$$

where $f_B \lesssim 1$ is a factor which controls the field strength in terms of the equipartition value, that is, $f_B = 1$ means the equipartition between the field and thermal energies. For fiducial values of $\Omega_{\rm b} = 0.06$ and $h = 0.75$, this gives for $f_B = 1$ the fields of 0.52-1.64 μgauss for $kT_{\rm obs} = 1{-}10$ keV at $z = 0$. It is much harder to estimate the fields in the unshocked infalling flow upstream to the shock, so here we will simply assume $B_1 = B_2(\rho_1/\rho_2)$ for a turbulent field. We note here that the turbulent amplification of the field could generate some fields not only downstream but also upstream to the shocks due to turbulent nature of the flows (Kulsrud et al. 1996).

We consider two kinds of models for the particle diffusion. The theoretical minimum for the diffusion coefficient in a strongly turbulent field with parallel geometry (i.e. the mean field is parallel to the flow direction) is given by the Bohm formula (see, e.g. , Drury 1983).

$$\kappa_{\rm B} = \frac{r_{\rm g}\, v}{3}, \tag{5}$$

where $r_{\rm g}$ is the gyroradius and v is the velocity of the particle. On the other hand, the minimum diffusion coefficient in the perpendicular shocks (i.e. the mean field is perpendicular to the flow direction) (Jokipii 1987) is given by

$$\kappa_{\rm J} = r_{\rm g}\, V_{\rm s} = 3\,(V_{\rm s}/v)\,\kappa_{\rm B}. \tag{6}$$

For high energy particles, $v \sim c$, so $\kappa_{\rm J}$ is smaller than $\kappa_{\rm B}$ by a factor of $V_{\rm s}/c$.

The rate of diffusive acceleration is determined by the diffusion coefficients $\kappa = \kappa_\parallel \cos^2\theta + \kappa_\perp \sin^2\theta$, where $\kappa_\parallel$ and $\kappa_\perp$ are the diffusion coefficient parallel and perpendicular to the magnetic field, respectively, and θ is the angle between the field and the shock normal (Jokipii 1987). Here we will follow Jokipii (1987) in assuming that $\kappa_\perp/\kappa_\parallel = (1+\eta^2)^{-1}$ and $\kappa_\parallel = \eta r_{\rm g} v/3$, where η is the ratio of the mean free path parallel to the magnetic field to the gyroradius. This is the result from standard kinetic theory which does not take account of field-line meandering. The Bohm diffusion coefficient corresponds to the case where the field is turbulent on a scale of $r_{\rm g}$ *at all momenta*, so it is effectively equivalent to the case of $\eta \to 1$ and $\kappa_\perp \sim \kappa_\parallel$ with a random distribution of the obliquity. On the other hand, the minimum diffusion coefficient in a perpendicular shock, which is referred as the Jokipii diffusion throughout this paper, corresponds to the case of $\theta = 90^\circ$ and $\eta_{\rm max} \sim c/(3V_s) = 100/V_{\rm s,3}$, where $V_{\rm s,3}$ is the shock velocity in units of 10^3 km s^{-1}. The latter is derived from the condition that $\kappa_{\perp,\rm min} \sim r_{\rm g} V_{\rm s} \sim (\eta r_{\rm g} c)/[3(1+\eta^2)]$, which implies that the particles should be scattered before they drift through the shock. This condition is necessary to maintain the isotropy of the particle distribution.

2.4 Energy Losses and Maximum Energy

The mean acceleration time scale for a particle to reach a momentum p is determined by the velocity jump at the shock, and the diffusion coefficient (e.g. , Drury 1983; Jokipii 1987), that is

$$\tau_{\rm acc} = \frac{\chi}{\chi - 1}\,\frac{v}{V_{\rm s}}\,\frac{\eta r_{\rm g}}{V_{\rm s}}\ \mathcal{J}(\chi,\theta) \tag{7}$$

where χ is the compression ratio of the shock and

$$\mathcal{J}(\chi,\theta) = \left[\cos^2\theta + \frac{\sin^2\theta}{1+\eta^2}\right] + \frac{\cos^2\theta + \chi^2\sin^2\theta/(1+\eta^2)}{\left[\cos^2\theta + \chi^2\sin^2\theta\right]^{3/2}}\ .$$

In the strong shock limit (i.e. $\chi = 4$), it can be written as

$$\tau_{\rm acc} = (4.23{\times}10^8\ {\rm years}) \times \frac{\eta E_{18}}{B_\mu V_{\rm s,3}^2}\ \mathcal{J}(4,\theta)\,, \tag{8}$$

where E_{18} is the particle energy in units of EeV, B_μ is the field strength in units of microgauss.

The protons lose energy due to pair production and pion production on the CBR not only on their way to earth, but also during their acceleration at the cluster shocks. At energies around 10 EeV the energy loss due to pair production is dominant and the loss time scale can be as short as $(5{\times}10^9{\rm years})/(1+z)^3$. Above the GZK cutoff at about 60 EeV photopion production becomes dominant and the loss time scale can be as short as $(5{\times}10^7{\rm years})/(1+z)^3$. The maximum energy $E_{\rm max}$ up to which the protons can be accelerated by the cluster accretion shocks is found by setting $\tau_{\rm acc} = \tau_{\rm int}$, with $\tau_{\rm int}$ being the time scale for interaction losses with the CBR, given by

$$\tau_{\rm int} = \frac{2\pi^2\hbar^3 c^2\gamma^2}{k\Theta(z)} \times \left[\int_{\epsilon_0'}^{\infty} d\epsilon'\,\epsilon'\,\langle\sigma\kappa\rangle(\epsilon')\ \ln\left[1 - \exp\left(\frac{-\epsilon'}{2\gamma k\Theta(z)}\right)\right]\right]^{-1}, \tag{9}$$

where γ is the Lorentz factor of the cosmic ray protons, $\Theta(z) = \Theta_0(1+z)$ the temperature of the cosmic microwave background at epoch z, and $\langle\sigma\kappa\rangle$ the inelasticity weighted cross section of the $p\gamma$ reaction, averaged over all final states (Rachen & Biermann, 1993). Equation (9) considers both pair production and pion production on the CBR; the impact of a putative cosmic infrared background is negligible for the determination of $E_{\rm max}$.

In Fig. 2 we have shown the loss time scales due to the interactions with CBR at $z = 0$, and the acceleration time scales in both Bohm and Jokipii diffusion limits. Here the acceleration time scales for $V_{\rm s} = 1000\,{\rm km\,s^{-1}}$and $B_1 = 1\mu$gauss are plotted. For other values of $V_{\rm s}$ and B, the acceleration time scales can be found by the following scaling relations: $\tau_{\rm acc,B} \propto V_{\rm s}^{-2}B^{-1}$ for the Bohm limit, and $\tau_{\rm acc,J} \propto V_{\rm s}^{-1}B^{-1}$ for the Jokipii limit. The maximum energy $E_{\rm max}$ for a shock of these parameters can be found by the intersection of two curves of $\tau_{\rm int}$ and $\tau_{\rm acc}$. The characteristic shape of loss time scale $\tau_{\rm int}$ as a function of energy causes only a weak dependence of $E_{\rm max}$ on $\tau_{\rm acc}$, if $\tau_{\rm acc} \lesssim 3{\times}10^9$ years, but

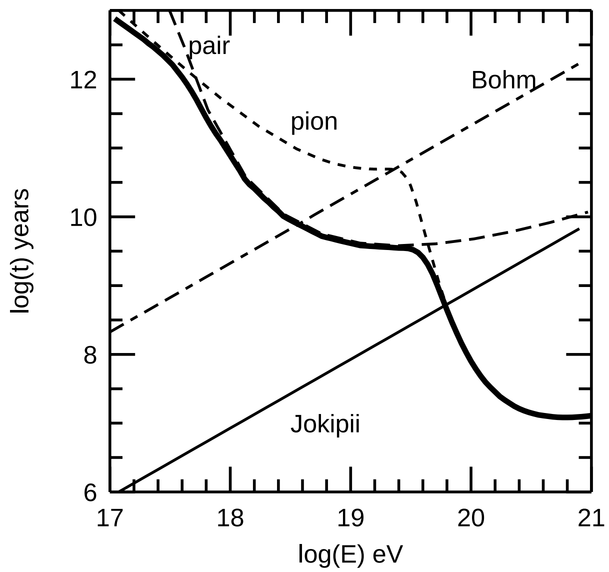

Figure 2. Schematic representation of acceleration and energy loss time scales as functions of particle energy. The dotted and dashed lines represent the loss time scale due to pion-production and pair-production, respectively, on both CMB and infrared background radiation. The heavy solid line is the added loss time scale due to both interactions, and it turns out that the infrared contribution is nowhere relevant. The dot-dashed and light solid line represent the acceleration time scales for Bohm and Jokipii diffusion models, respectively. For this case, adopted shock parameters are $V_{\rm s} = 1000\,{\rm km\,s^{-1}}$and $B = 1\mu$gauss. The intersection point of both curves determines the maximum energy of acceleration.

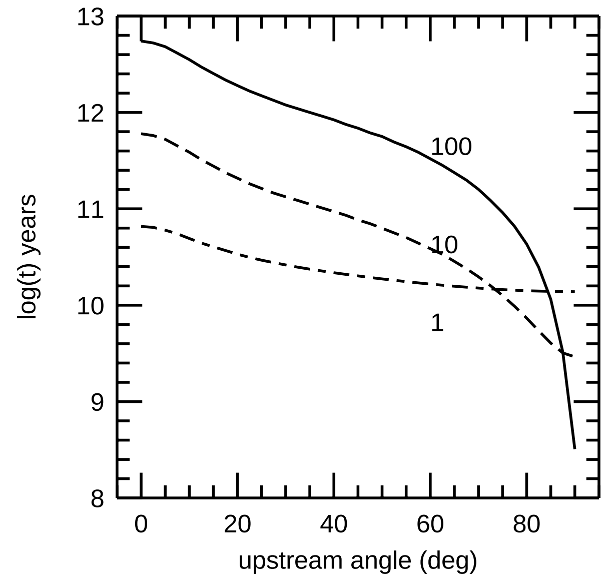

Figure 3. Time scales for the protons to be accelerated to $E = 10^{19.6}$eV by a shock of $V_{\rm s} = 1000\,{\rm km\,s^{-1}}$and $B_1 = 1\mu$gauss as a function for obliquity. The curves are labeled with the values of $\eta = \lambda/r_{\rm g} = 1$, 10, and 100.

a strong dependence if the acceleration is slower. This time scale corresponds to $\tau_{\rm int}$ for $E \sim 10^{19.6}$eV, Thus one can see from Fig. 2 that, for a canonical accretion shock, the maximum energy for the Bohm limit is most likely to be set by the pair-production loss and ranges from 1 to 10 EeV, while that for the Jokipii limit is set by the pion-production loss at around 50 EeV.

According to equation (8), the acceleration time scale is in general dependent upon the obliquity and the strength of turbulent field. Fig. 3 shows the acceleration time scales for the particles to reach up to $E = 10^{19.6}$ eV for $\eta = 1, 10$ and 100 as a function of the upstream oblique angle θ. Once again the acceleration time scales for $V_{\rm s} = 1000\,{\rm km\,s^{-1}}$and for $B_1 = 1\mu$gauss are plotted. This shows that the particles can be accelerated above $10^{19.6}$eV only when the mean field is nearly perpendicular to the shock normal. The obliquity dependence of $\tau_{\rm acc}$ is rather weak in the strong scattering limit ($\eta \sim 1$), while it becomes very strong when the cross-field diffusion is small ($\eta \sim 100$). One can also note that $\tau_{\rm acc}$ linearly increases with η in quasi-parallel shocks (i.e. $\theta \approx 0^\circ$), so the acceleration is most efficient in the limit of strong turbulences (i.e. $\eta \rightarrow 1$). On the other hand, $\tau_{\rm acc}$ is nearly inversely proportional to η in quasi-perpendicular shocks (i.e. $\theta \approx 90^\circ$), so the acceleration becomes most efficient in the limit of weak scattering (i.e. $\eta \rightarrow \eta_{\rm max}$).

2.5 Injection Spectrum

For each cluster with given values of $kT_{\rm obs}$ and z, the self-similar solution gives the shock parameters, ρ_1, $V_{\rm s}$, $r_{\rm s}$, and B_1, and the time scale condition in equation (9) gives the estimate for the cut-off energy in the particle spectrum. Momentum distribution of the protons at the shock is assumed to be a power law whose index is given by the parameter α, and with an exponential cutoff at the maximum energy ($E_{\rm max} = cp_{\rm c}$). The minimum momentum for the power law is $p_{\rm inj} \sim m_p V_{\rm s}$. The power-law index is $\alpha = 4$ in the limit of strong shocks with dynamically insignificant CR energy density (i.e. test-particle limit), but can vary slightly around this value for real shock waves (we note that $\alpha = 4$ corresponds to a power-law in the energy spectrum of E^{-2}). Here we consider values of α equal to and slightly larger than 4. We further assume that the process with which the particles escape from the shock is momentum-independent, that is, the particle spectra injected into intergalactic (IG) space has the same shape of the proton spectra at the shock. Then the injected distribution from a cluster with a given ICM temperature is given by

$$f(T,z,p) = A(T,z)\, p^{-\alpha} \exp\left[-\frac{p}{p_{\rm c}(T,z)}\right] \quad , \tag{10}$$

where $A(T,z)$ is a normalization constant. It is chosen by assuming that a small fraction (ϵ) of the kinetic energy density of the infalling matter in the shock rest frame ($\rho_1 V_{\rm s}^2$) is converted to CR energy and then injected into IG space. So the free parameter, ϵ, along with the assumed value of α controls the amplitude $A(T,z)$ according to

$$\epsilon = \frac{4\pi}{\rho_1 V_{\rm s}^2} \int_{p_{\rm inj}}^{p_{\rm c}} f(T,z,p)\, E\, p^2\, dp \; . \tag{11}$$

Thus, for given values of ϵ and α, the normalization constants were numerically calculated for given values of T and z. For $\alpha = 4$, this gives an approximation for $A(T, z)$ such as

$$A(T,z) \simeq \frac{\epsilon \rho_1 V_{\rm s}^2}{4\pi m_p c^2 \ln{(p_{\rm c}/m_p c)}} \quad . \tag{12}$$

This along with $p_{\rm c} = (1/c)\, E_{\rm max}(T, z)$ completely specifies the injection spectrum from a cluster.

In KRJ96, it was assumed that the particles escape from the shock once they diffuse upstream to a distance comparable to the radius of the shock due to the lateral diffusion. Since this escape process depends on the diffusion length, the particle spectrum injected into the IG space has a strong dependence on the momentum, that is, only highest energy particles with longest diffusion length can diffuse out from the shock. We note that, if we take the same momentum-dependent escape model, the resulting spectrum cannot fit the observed CR distribution over a wide range of the particle energy as well as the spectrum of our momentum-independent escape model can. We will also discuss the escape process in §3.1.

2.6 Cluster Distribution Function

The evolution of the cluster population is a critical issue which has been under intense discussion recently, since it can provide a potentially powerful tool for discriminating different cosmological models, especially the matter density of the universe (i.e. Ω_0). Theoretical prediction based on a hierarchical clustering model is that clusters are fainter but more abundant in the past (Kaiser 1986). The evolution of cluster population can be calculated by using various methods based on so-called Press-Schechter formalism, if the background cosmology (i.e. Ω_0 and Λ_0) and the initial density power spectrum are given (e.g. Eke, Cole & Frenk 1996; Kitayama & Suto 1996; Bond & Myers 1996). The density fluctuations of cluster mass ($M \sim 10^{15} M_\odot$) would form earlier in low Ω_0 universe than in high Ω_0 universe, while their formation rate at the present epoch would be higher in low bias models (i.e. larger σ_8, Cen & Ostriker 1994; Kitayama & Suto 1996; Bond & Myers 1996). Theses studies as well as most numerical studies (Kang et al. 1994a; Tsai & Buote 1996) indicate that the standard CDM model of a critical density universe produces too many clusters compared to observed local cluster abundance. Also recent observations of distant X-ray clusters at $z < 0.3$. (Castander et al. 1995; Ebeling et al. 1995) seem to indicate that clusters have evolved very little or not evolved at all in this redshift range, which is more consistent with low bias or low density models.

In order to model the distribution of cosmological population of clusters, first we adopted a temperature distribution function of clusters given by Henry and Arnaud (1991) which was derived from the observed local clusters. The number density in unit comoving volume at the present is given by for kT =1–10 keV,

$$n_{\rm o}(kT) = 1.8{\times}10^{-3}\, h^3\, {\rm Mpc}^{-3}\, {\rm keV}^{-1} \left(\frac{kT}{\rm keV}\right)^{-4.7} \tag{13}$$

Secondly, we made a simple assumption that the comoving density of cluster evolves as a power-law of a scale factor, $(1+z)$, that is, $n(kT, z) = n_{\rm o}(kT)(1+z)^m$. The value of m at redshifts $z < 0.3$ is most likely zero or slightly negative according to the observations mentioned above. According to the scaling law of Kaiser (1986) for a scale free power spectrum in a critical density universe, for CDM like density power spectrum of $n_{\rm eff} \sim -1$ at the cluster mass scale, the index m would be -0.7 for the above power-law temperature distribution given by equation (13). Here we will consider a range of values for m ($-1 \le m \le +1$). We take the minimum value of the redshift as that of the Virgo cluster, that is, $z_{\rm min} = 0.0036$. Most X-ray bright, rich clusters would form after $z \sim 5$ and in fact clusters of $M = 10^{15} M_\odot$ form around $z_{\rm f} \sim 0.6$ in an $\Omega_0 = 1$ universe. From these considerations, we set the maximum redshift to be $z_{\rm max} = 5$. But in fact contributions from the epoch earlier than $z \sim 1$ are insignificant even for $m = 1$ modeled because of the spatial dilution ($1/r^2$ factor) and the interactions with CBR for distant sources.

Then number of clusters between z and $z + dz$ for a given temperature kT is given by

$$dN(kT, z) = n_{\rm o}(kT)\; (1+z)^m \left[\frac{dV_{\rm c}}{dz}\right]\; dz \tag{14}$$

where $dV_{\rm c}/dz = 4\pi(c/H_0)^3[1-(1+z)^{-1/2}]^2/(1+z)^{3/2}$ for $\Omega_0 = 1$ and $(dV_{\rm c}/dz)\, dz$ is the comoving volume of a shell defined between z and $z + dz$. Adding up the contribution from clusters from $z_{\rm min}$ to $z_{\rm max}$ while including the interactions with CBR, and integrating over the temperature distribution, one can get the particle flux observed at earth according to

$$\begin{aligned} J(E) \quad &= \quad \frac{c}{4} \int_{kT_{\rm min}}^{kT_{\rm max}} d(kT) \int_{z_{\rm min}}^{z_{\rm max}} dz \\ &\left\{ \left[\frac{dN(kT,z)}{dz}\right] \left[\frac{r_{\rm s}}{d_{\rm cl}(z)}\right]^2 M(E)\, f(T,z,E) \right\} . \end{aligned} \tag{15}$$

$d_{\rm cl}(z) = (6\, h^{-1}{\rm Gpc})\,(1+z-\sqrt{1+z})$ is the luminosity distance of a cluster at redshift z, given in an $\Omega_0 = 1$ cosmology. $M(E)$ is the modification factor that accounts for the interactions with CBR along the pathway to earth, calculated as defined in Rachen and Biermann (1993), using the the continuous energy loss approximation introduced by Berezinsky and Grigor'eva (1988). It has been pointed out that the characteristic spikes occurring in $M(E)$ close to the cutoff are exaggerated because of the continuous loss approximation (Yoshida & Teshima 1993; Protheroe & Johnson 1996). In the integration over redshift, however, those features are smoothed out, and it should only be noted that the cutoff is not quite as sharp as proposed by this method.

3 RESULTS AND DISCUSSION

Fig. 4 shows the maximum energy calculated according to equation $\tau_{\rm acc} = \tau_{\rm int}$ for both diffusion models for clusters with kT = 2, 4, 6, and 10 keV at $z = 0.0$ as a function of the parameter f_B. As found in KRJ96, the maximum energy cannot go above the GZK-cutoff for the Bohm diffusion coefficient. Thus if the general field direction is radial in the accretion flow, then the diffusive acceleration is too slow to accelerate the protons above the GZK-cutoff. We note from Fig. 3, however, that the transition from quasi-perpendicular Jokipii diffusion to Bohm diffusion does not lead to a linear decrease of $E_{\rm max}$, so that the model can work also for moderately oblique shocks with $\eta \gg 1$.

Fig. 5 shows the expected proton spectrum from the cluster accretion shocks calculated for both diffusion models. The power-law index for the cosmological evolution is $m = 0$, the magnetic field strength parameter is $f_B = 1$ for all models. Different values of the injection energy fraction ϵ is assumed for each value of the spectral index α to obtain the better fit with the observation around 10 EeV for Jokipii models and around 1 EeV for Bohm models, respectively. As expected, higher ϵ is required for steeper spectra. For Jokipii diffusion, the models with $4.0 \le \alpha \le 4.2$ show good fits

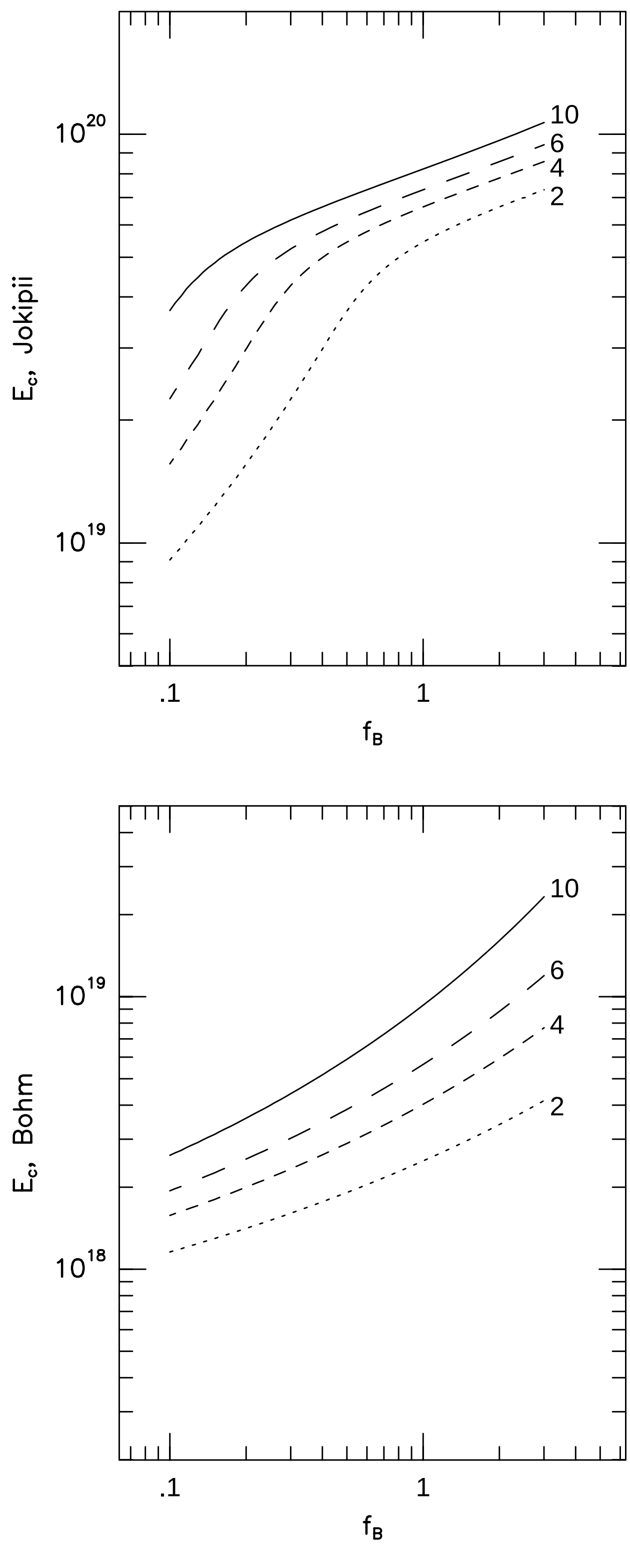

Figure 4. Maximum energy of the protons accelerated by the accretion shocks associated with the clusters with ICM temperatures of $kT = 2,\ 4,\ 6,$ and 10 keV at $z = 0$ as a function of the parameter f_B ($\Omega_\mathrm{b} = 0.06$ and $h = 0.75$). The top panel shows models with Jokipii diffusion, the bottom panel shows models with Bohm diffusion.

to the data, while the models with Bohm diffusion produce too few particles above 1 EeV.

The sensitivity to the cluster evolutionary model is shown in the top panel of Fig. 6. For all models here $\alpha = 4, \epsilon = 5{\times}10^{-5}$, and $f_B = 1$. The power-law index for the evolution with the redshift for the comoving density of clusters are $m = -1, 0,$ and $+1$ for the dotted, solid, and dashed lines, respectively. Since the nearby clusters contribute most, the resulting spectrum is not severely dependent on the evolutionary model at high redshifts. The relatively weak cosmological evolution of galaxy clusters compared to radio galaxies provides an even better fit to the light component data derived from the Fly's Eye air shower analysis (Rachen et al. 1993), but we point out that the errors are large here and that these data have not been confirmed by the AGASA collaboration.

The bottom panel of Fig. 6 shows how the results depend on the magnetic field strength. For all models here $\alpha = 4,\ \ \epsilon = 5{\times}10^{-5}$, and $m = 0$ is used. The field strength relative to the equipartition strength is represented by $f_B = 0.1, 0.5$ and 1 for the dashed, dotted, and solid lines. One can see that cluster accretion shock cannot produce enough particles above 10 EeV, if the field strength is much smaller than the equipartition value (i.e. $f_B \lesssim 0.1$). The reduction of the field from the equipartition by a factor of up to two, however, still can give a reasonable fit to the observations.

3.1 The escape process

We need to consider why nature would choose an injection efficiency at a level of order 10^{-4}. It is believed that a supernova blast wave can transfer a few $(1{-}10)$% of the explosion energy into the cosmic ray component and the particles accelerated by the remnant shock are injected into the interstellar medium when the shock either decays into a sonic wave or breaks up due to various instabilities. It would be natural to assume that a similar fraction of the kinetic energy of the infalling flow is likely to be transferred to the CRs. Unlike a supernova blast wave which expands from a point-like explosion into the interstellar space, however, an accretion shock forms due to the converging flow associated with gravitational collapse. The particles will be mostly confined near the shock and so the escape of particles by swimming against infalling flows upstream will be very rare. This calls for some physical models of escape of CRs from the cluster shocks which can explain that only a few percent of accelerated particles (which contains only a few percent of the infall energy) is injected into the intergalactic space. Here we must note that the escape process cannot be energy dependent, because then the spectrum could not match the observations. In case of galactic cosmic rays which travel against the solar wind to reach the earth, strong attenuation is severely energy dependent, and so the solar wind does not provide a good example apparently.

For clusters which are connected with filaments or supergalactic planes, the particles can escape from the clusters along the filaments/sheets since the field lines are mostly likely parallel along those structures according to the flow velocity inside the filaments/sheets (see, for example, Fig. 1 of KRJ96). In such a picture the low apparent efficiency is simply a measure of the difficulty of feeding the particles into the sheets from the surrounding of a cluster. If this conjecture is right, the arrival direction of UHECRs should show a correlation with the Supergalactic plane, and so identifying a source cluster with a particular CR event may become less stringent. The particles can also leak when the accretion shock becomes unstable and breaks up temporarily due to strong activities

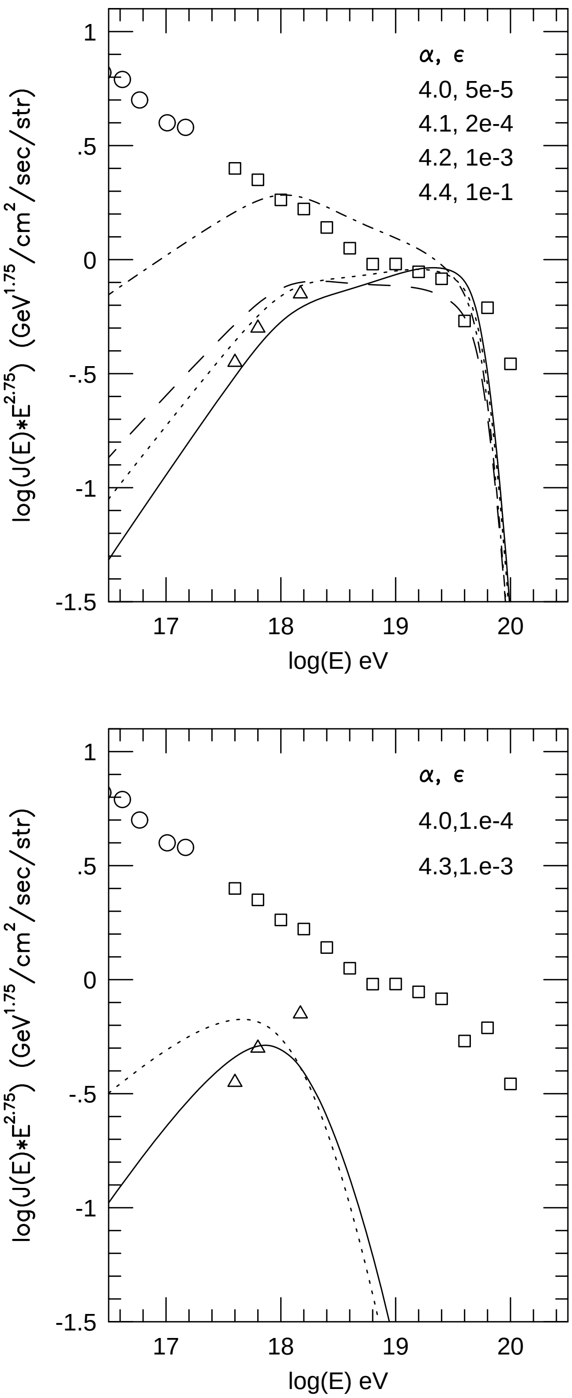

Figure 5. Estimated proton flux from a cosmological ensemble of cluster accretion shocks. The top panel shows the models with Jokipii diffusion. The solid line is for $\alpha = 4.0, \epsilon = 5\times10^{-5}$, dotted line for $\alpha = 4.1$, $\epsilon = 2\times10^{-4}$ dashed line for $\alpha = 4.2$, $\epsilon = 10^{-3}$ and dot-dashed line for $\alpha = 4.4$, $\epsilon = 10^{-1}$. The bottom panel shows the results for Bohm diffusion. The solid line is for $\alpha = 4.0$, $\epsilon = 10^{-4}$, and dotted line for $\alpha = 4.3$, $\epsilon = 10^{-3}$. In all models we set $f_B = 1$, $m = 0$, $\Omega_{\rm b} = 0.06$ and $h = 0.75$. The data points represent the all-experiment data set collected by T. Stanev (squares, see Rachen, Stanev & Biermann 1993), and the light component inferred from the Fly's Eye composition measurements (triangles, ibid).

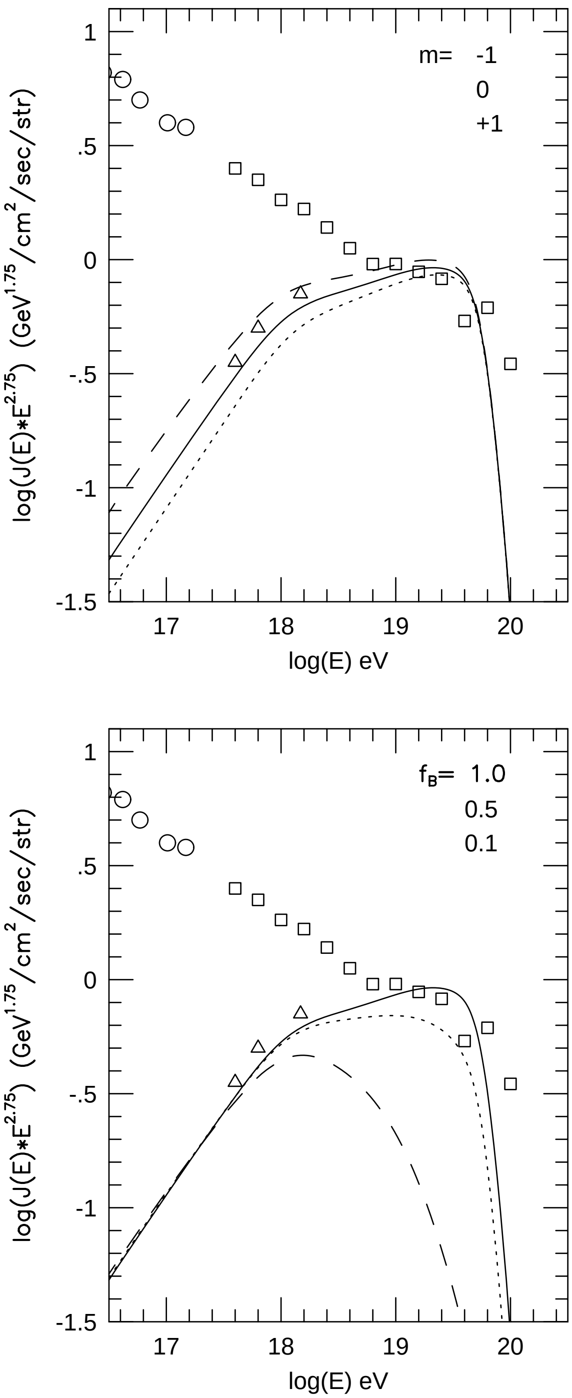

Figure 6. Estimated proton flux from a cosmological ensemble of cluster accretion shocks. For all models here Jokipii diffusion is assumed and $\epsilon = 5\times10^{-5}$, $\alpha = 4.0$, $\Omega_{\rm b} = 0.06$ and $h = 0.75$. The top panel shows the models with $f_B = 1$, but with different values of $m = -1, 0, +1$ (dotted, solid, and dashed lines, respectively). The bottom panel shows the models with $m = 0$, but with different values of $f_B = 0.1, 0.5, 1.0$ (dashed, dotted, and solid lines respectively). The data points are the same as in Fig. 5.

of AGNs inside the cluster. This could be the only possible escape route for the particles accelerated by isolated clusters which are not part of superclusters. The issue is closely related with the topology and the turbulent nature of the magnetic fields in the sheets, filaments and around clusters and will be discussed in upcoming papers.

3.2 The magnetic fields outside clusters

For our model to work, the magnetic field just outside clusters of galaxies must be of order 0.1 microgauss, but cannot be less, since then the Larmor radius condition would fail, i.e. the particles could not be contained. Does this imply that we require a primordial magnetic field?

We note that in the evolution of clusters, repeated formation of giant radio galaxies is likely, and they can break out and pollute the environment over very large distances with magnetic fields. Therefore, with such a model tracing the large scale magnetic fields to very large radio galaxies, one would not need a primordial magnetic field.

Of course, a primordial magnetic field would have to be weaker than the observational constraints, given by the cosmological limits on the Faraday rotation of distant radio sources; a primordial magnetic field is likely nearly homogeneous inside the bubble structure of the galaxy distribution, and so an upper limit of ≈ 200 picogauss would obtain, with a substantial strengthening towards the spatial environment near to clusters. If we use an inverse Parker wind model (as an analogy to the expansion of the gas from a radio galaxy) to do such an extrapolation, i.e. scale the magnetic field strength inversely with the distance to a cluster, then such an upper limit is just about compatible with the magnetic fields required by our model.

In conclusion, we cannot decide between the two possibilities, a primordial magnetic field, and one that derives from radio galaxies, or some other sources such as normal galaxies, even on the large scale.

3.3 Cosmic rays inside clusters

The accretion shock produces cosmic rays scaling with the overall energetics of the cluster, but with only some fraction of the shock kinetic energy; thus the energy density of cosmic rays inside clusters is expected to be less than 10% of the thermal energy content in our model of the intracluster medium. On the other hand, the cosmic rays injected from radio galaxies may produce magnetic field strengths and energy densities of cosmic rays close to equipartition: Enßlin et al. (1996) proposed that radio galaxies can push the cosmic ray content in clusters to the stability limit such as is believed to happen in the interstellar medium due to supernova explosions (see, e.g. Parker 1969 for a review). As demonstrated by Enßlin et al. (1996) these two alternatives may be decidable in the near future with gamma-ray observations due to the pion decay production in pp-collisions.

3.4 Arrival directions of UHECRs

There is a study of the Haverah Park events by Stanev et al. (1995), which shows that the arrival directions of the UHECRs are not uniform, but seem to show positive correlations with the supergalactic plane; a related study of the AGASA events by Hayashida et al. (1996) supports a connection between the supergalactic plane and the arrival directions of at least a significant fraction of UHE CRs. This supports models for the extragalactic origin of the highest energy cosmic rays, especially the ones in which the sources are likely to be associated with the supergalactic plane. They include radio galaxies and cluster accretion shocks. We have looked at the distribution of local clusters ($z < 0.3$) and generated a map of the cluster accretion shocks weighted by the particle flux above a few times 10 EeV (Kang, Rachen & Biermann 1996). The map shows the similar degree of correlations with the supergalactic plane as the arrival directions of the UHECRs found in Stanev et al. (1995) and Hayashida et al. (1996). This is consistent with the fact that rich clusters beyond the local supergalactic plane ($d > 30h^{-1}$Mpc) such as the Coma (8.3keV), Perseus (6.3keV), 3C129 (6.2keV), A3571 (7.6keV), and Centaurus (3.9 keV) clusters are at low supergalactic latitudes. These rich clusters are dominant contributors of UHE CRs according to our model. The Virgo cluster is the most prominent cluster in the Local Supercluster, but its temperature is only only 2.4 keV. Thus its contribution through the accretion shock is important only for energies below GZK cutoff; of course, it may have an additional contribution from its radio galaxies.

3.5 The most energetic events

Concerning the events above 100 EeV, the present model can give a straightforward explanation for the enhancement of highest energy events detected by the Haverah Park experiment at high north galactic latitudes (see Stanev et al. 1995, and references therein), since the Virgo and Coma clusters are in that direction. On the other hand, the applicability of the hot spot acceleration model to the center-brightened radio galaxies M87 harbored by the Virgo cluster is not yet clear. Whether the 320 EeV Fly's Eye event may be explained in the context of cluster accretion shock acceleration, remains to be investigated. The AWM7 cluster ($T = 4.0$ keV, $z = 0.0176$) might be the closest rich cluster to the general direction of this event. Alternatively, there is a candidate radio galaxy, 3C134; its redshift is unknown due to obscuration, but its distance can be estimated from radio size-luminosity relations to be in the range $30-300$ Mpc, thus it may well be the closest FR-II radio galaxy (Rachen 1995). This estimate has been recently confirmed by an infrared detection of the central galaxy (Hartmann 1996). All in all, since the maximum energy of acceleration may vary strongly between individual objects in both models, one may have to consider both possibilities in the explanation of the highest energy events.

4 CONCLUSIONS AND OUTLOOK

The studies on the expected CR spectrum and distribution from galaxy cluster accretion shocks seem to give strong evidences for a significant UHECR contribution, which could reach up to the highest observed energies. Some aspects of the model need further detailed considerations, however, since it has to rely on several somewhat speculative assumptions. A first question could be if one can prove observationally the existence of large-scale accretion shocks around clusters. Although the shocks and the hot component of IGM can be clearly seen to exist in most of cosmological hydrodynamic simulations, independent of the details of cosmological models; it is observationally challenging to detect such shocks because the hot gas does not radiate or absorb photons at a detectable level. But the hot gas of $T = 10^5-10^6$ K heated by the large-scale

accretion shocks seems to be a major component of the intergalactic medium in the simulated universes (Nath & Biermann 1993; Ostriker & Cen 1996). One can also ask if the model would work for lower density universes such as open CDM models and flat CDM models with cosmological constant, because it is expected that the accretion flows are weaker in such cosmologies.

Another crucial assumption for our model is the magnetic field in equipartition with the thermal energy in the vicinity of accretion shocks. The general topology and turbulent nature of the large-scale magnetic field are also closely related with our particle diffusion model and the escape process. We note that the Jokipii expression for the possible range of the diffusion coefficient in an oblique shock geometry has been derived going to the absolute limit of what is geometrically possible; it may be an extreme overstatement of the limit relevant in nature. Therefore, we emphasize that we use this as an assumption going into the modelling. On the other hand, considering the loss time scale and comparing the two limits, Bohm and Jokipii (Fig. 2), the maximum energy possible is only moderately changed, even when we are 1/3 of the way between Jokipii and Bohm limits. Therefore, we conclude that the speculative nature of the Jokipii *limit* has relatively little influence on our essential result, as long as nature allows us to go two thirds of the way towards Jokipii's limit. This demonstrates that the theoretical studies on how a magnetic field is generated and amplified during the structure and galaxy formation (e.g. Kulsrud et al. 1996; Ryu, Kang & Biermann 1996; Biermann 1996), as well as the observational studies on how the real magnetic field in the IGM is distributed are very important to further development of our model.

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